Dominant $2\pi\gamma$-exchange nucleon-nucleon interaction: Spin-spin and tensor potentials

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Abstract

We calculate at two-loop order in chiral perturbation theory the electromagnetic corrections to the two-pion exchange nucleon-nucleon interaction that is generated by the isovector spin-flip $\pi\pi NN$ contact-vertex proportional to the large low-energy constant $c_4 \simeq 3.4\,\mathrm{GeV}^{-1}$. We find that the respective $2\pi\gamma$-exchange potentials contain sizeable isospin-breaking components which reach up to -4% of corresponding isovector 2π-exchange potentials. The typical values of these novel charge-independence breaking spin-spin and tensor potentials are $-0.11\,\mathrm{MeV}$ and $0.09\,\mathrm{MeV}$, at a nucleon distance of $r = m_\pi^{-1} = 1.4\,\mathrm{fm}$. The charge-symmetry breaking spin-spin and tensor potentials come out a factor of 2.4 smaller. Our analytical results for these presumably dominant isospin-violating spin-spin and tensor NN-forces are in a form such that they can be easily implemented into phase-shift analyses and few-body calculations.

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Isospin-violation in the nuclear force is a subject of current interest. Significant advances in the understanding of nuclear isospin-violation have been made in the past years by employing methods of effective field theory (in particular chiral perturbation theory). Van Kolck et al. [1] were the first to calculate (in a manifestly gauge-invariant way) the complete leading-order pion-photon exchange nucleon-nucleon interaction. In addition, the charge-independence and charge-symmetry breaking effects arising from the pion mass difference $m_{\pi^+} - m_{\pi^0} = 4.59\,\mathrm{MeV}$ and the nucleon mass difference $M_n - M_p = 1.29\,\mathrm{MeV}$ on the (leading order) two-pion exchange NN-potential have been worked out in Refs. [2, 3]. Recently, Epelbaum et al. [4] have continued this line of approach by deriving the subleading isospin-breaking 2π-exchange NN-potentials and classifying the relevant isospin-breaking four-nucleon contact terms.

In a recent work [5] we have calculated the electromagnetic (i.e. one-photon exchange) corrections to the dominant two-pion exchange NN-interaction. The latter comes in form of a strongly attractive isoscalar central potential generated by a one-loop triangle diagram involving the $\pi\pi NN$ contact-vertex proportional to the large low-energy constant $c_3 \simeq -3.3\,\mathrm{GeV}^{-1}$. The dynamics behind this large value of c_3 is the excitation of the low-lying $\Delta(1232)$-resonance. It has been found in ref. [5] that this particular class of two-loop $2\pi\gamma$-exchange diagrams (proportional to c_3) leads to sizeable charge-independence and charge-symmetry breaking central NN-potentials which amount to 0.3 MeV at a nucleon distance of $r = m_\pi^{-1} = 1.4\,\mathrm{fm}$. These are in fact the largest isospin-violating NN-potentials derived so far from long-range pion-exchange [6]. The effects from the other equally large low-energy constant $c_2 \simeq 3.3\,\mathrm{GeV}^{-1}$ have also been studied in ref. [5]. Isospin-violating potentials, $\widetilde{V}_C^{(\mathrm{cib})}(m_\pi^{-1}) = -23.4\,\mathrm{keV}$ and $V_C^{(\mathrm{csb})}(m_\pi^{-1}) = -14.8\,\mathrm{keV}$, more than an order of magnitude smaller have only been found. A reason for this suppression of the effects from the c_2-vertex is that the relevant two-loop spectral function stems in this case from a $2\pi\gamma$ three-body phase space integral whose integrand vanishes on the phase space boundary.

The remaining large low-energy constant $c_4 \simeq 3.4\,\mathrm{GeV}^{-1}$ enters in an isovector spin-flip $\pi\pi NN$ contact vertex. When inserted into the two-loop $2\pi\gamma$-exchange diagrams the c_4-vertex generates isospin-violating spin-spin and tensor potentials. It is the purpose of the present short paper to calculate these remaining contributions to the isospin-violating NN-potential along the same lines as in ref.[5] (i.e. by computing analytically the two-loop spectral functions). We do indeed find sizeable charge-independence breaking spin-spin and tensor potentials with values of $-0.11\,\mathrm{MeV}$ and $0.09\,\mathrm{MeV}$, at a nucleon distance of $r = m_\pi^{-1} = 1.4\,\mathrm{fm}$. These novel charge-independence breaking spin-spin and tensor potentials are actually more than an order of magnitude larger than their counterparts arising from (lowest order) pion-photon exchange (see Table I in ref.[6]).

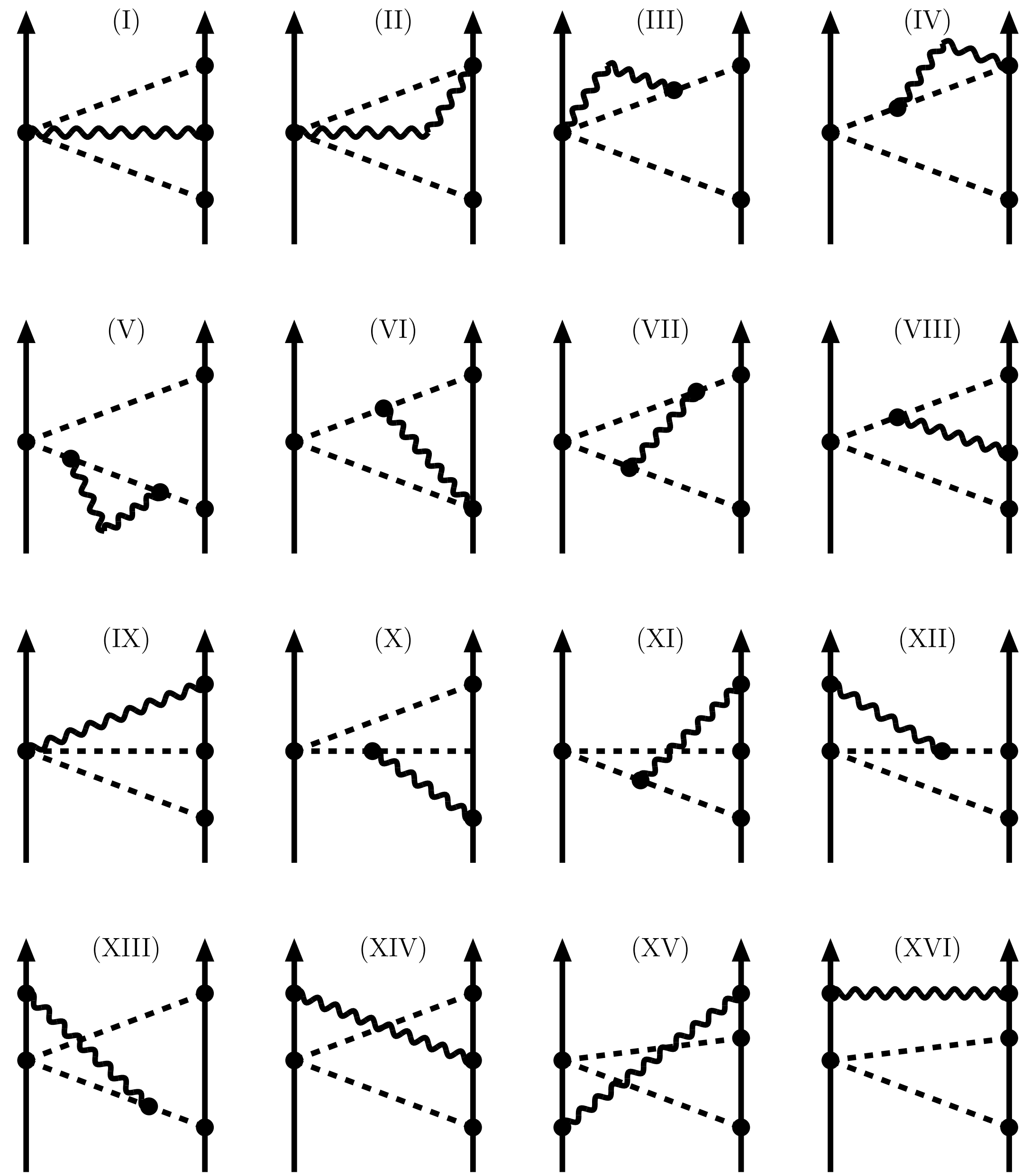

Fig. 1: Electromagnetic corrections to the dominant isovector spin-spin and tensor 2π-exchange NN-interaction generated by the $\pi\pi NN$ contact-vertex $ic_4 f_\pi^{-2}\,\epsilon^{abc}\tau^c\,\vec\sigma\cdot(\vec q_a \times \vec q_b)$. Diagrams with the contact-vertex at the right nucleon line and diagrams turned upside-down are not shown. The spectral functions $\mathrm{Im}V_{S,T}(i\mu)$ are calculated by cutting the intermediate $\pi\pi\gamma$ three-particle state.

In order to test their phenomenological relevance these novel and exceptionally large isospin-violating spin-spin and tensor potentials should be implemented into NN phase-shift analyses [7, 8, 9] and few-body calculations.

Let us start with recalling the chiral 2π-exchange NN-interaction. Its dominant contribution to the isovector spin-spin and tensor channels comes from a triangle diagram involving one $\pi\pi NN$ contact-vertex with momentum-dependent coupling: $ic_4 f_\pi^{-2}\, \epsilon^{abc}\tau^c\, \vec\sigma\cdot(\vec q_a\times\vec q_b)$. The corresponding potentials in coordinate space (obtained by a modified Laplace transformation) read [10]:

$$\widetilde W_S^{(2\pi)}(r) = \frac{c_4 g_A^2}{48\pi^2 f_\pi^4}\,\frac{e^{-2z}}{r^6}\,(1+z)(3+3z+2z^2)\,, \tag{1}$$

$$\widetilde W_T^{(2\pi)}(r) = -\frac{c_4 g_A^2}{48\pi^2 f_\pi^4}\,\frac{e^{-2z}}{r^6}\,(1+z)(3+3z+z^2)\,, \tag{2}$$

where $z = m_\pi r$. The occurring parameters are: $g_A = 1.3$ (the nucleon axial-vector coupling constant), $f_\pi = 92.4\,\text{MeV}$ (the pion decay constant), $m_\pi = 139.57\,\text{MeV}$ (the charged pion mass) and the low-energy constant $c_4 = 3.4\,\text{GeV}^{-1}$. The latter value has obtained in one-loop chiral perturbation theory analyses of low-energy πN-scattering [11] and it has also been used (successfully) by Epelbaum et al. [12] in calculations of NN-phase shifts at next-to-next-to-next-to-leading order in the chiral expansion. The numerical values in the first row of Tables I and II display the magnitude and r-dependence of these dominant isovector spin-spin and tensor 2π-exchange potentials $\widetilde W_{S,T}^{(2\pi)}(r)$. It should also be noted that $\widetilde W_T^{(2\pi)}(r)$ reduces at intermediate distances, $1\,\text{fm}\le r\le 2\,\text{fm}$, the too strong isovector tensor potential from 1π-exchange (see e.g. Fig. 9 in ref.[10]).

Now we will add to the 2π-exchange triangle diagram a photon line which runs from one side to the other. There are five positions for the photon to start on the left hand side and seven positions to arrive at the right hand side. Leaving out those four diagrams which vanish in Feynman gauge (with photon propagator proportional to $g_{\mu\nu}$) we get the 16 representative diagrams shown in Fig. 1. Except for diagram (I) these are to be understood as being duplicated by horizontally reflected partners. A further doubling of the number of diagrams comes from interchanging the role of both nucleons. The two-loop diagrams in Fig. 1 with the c_4-contact vertex included generate a contribution to the T-matrix of elastic NN-scattering of the form:

$$\mathcal{T}_{NN} = V_S(q)\ \vec\sigma_1\cdot\vec\sigma_2 + V_T(q)\ \vec\sigma_1\cdot\vec q\ \vec\sigma_2\cdot\vec q\,, \tag{3}$$

with $\vec\sigma_{1,2}$ the usual spin-operators and $\vec q$ the momentum transfer between the two nucleons. We are interested here only in the long-range parts of the associated coordinate-space potentials (disregarding zero-range $\delta^3(\vec r)$-terms). For that purpose it is sufficient to know the spectral functions or imaginary parts of the two-loop diagrams. Making use of (perturbative) unitarity in the form of the Cutkosky cutting rule we can calculate the two-loop spectral functions as integrals of the (subthreshold) $\bar NN\to\pi\pi\gamma\to\bar NN$ transition amplitudes over the Lorentz-invariant $2\pi\gamma$ three-particle phase space. In the (conveniently chosen) center-of-mass frame this leads to two angular integrations and two integrals over the pion energies. Due to the heavy nucleon limit ($M_N\to\infty$) and the masslessness of the photon ($m_\gamma = 0$) several simplifications occur and therefore most of these integrations can actually be performed in closed analytical form. For a concise presentation of our results it is furthermore advantageous to consider two particular linear combinations of the spin-spin and tensor spectral functions $\text{Im}V_{S,T}(i\mu)$ and to scale out all common (dimensionful) parameters:

$$3\,\text{Im}V_S(i\mu) - \mu^2\,\text{Im}V_T(i\mu) = \frac{\alpha c_4 g_A^2 m_\pi^3}{\pi(4f_\pi)^4}\ S_1\Big(\frac{\mu}{m_\pi}\Big)\,, \tag{4}$$

$$\mathrm{Im}V_S(i\mu) - \mu^2\,\mathrm{Im}V_T(i\mu) = \frac{\alpha c_4 g_A^2 m_\pi^3}{\pi(4f_\pi)^4}\; S_2\Big(\frac{\mu}{m_\pi}\Big)\,. \tag{5}$$

Here, $\mu \geq 2m_\pi$ denotes the $\pi\pi\gamma$-invariant mass and $\alpha = 1/137.036$ is the fine-structure constant. Without going into further technical details, we enumerate now the contributions of the 16 representative $2\pi\gamma$-exchange diagrams shown in Fig. 1 to the dimensionless spectral functions $S_1(u)$ and $S_2(u)$ which depend only on the dimensionless variable $u = \mu/m_\pi \geq 2$. We find three vanishing contributions:

$$S_1(u)^{(\mathrm{I})} = S_1(u)^{(\mathrm{VI})} = S_1(u)^{(\mathrm{IX})} = 0\,. \tag{6}$$

In the Feynman gauge this property is obvious for diagrams (I) and (IX), whereas one encounters for diagram (VI) a quite nontrivial double integral over the pion energies which surprisingly turns out to be equal to zero. The remaining nonvanishing contributions to the spectral function $S_1(u)$ read for $u \geq 2$:

$$S_1(u)^{(\mathrm{II})} = \vec\tau_1\cdot\vec\tau_2\left\{16 - 8u + \frac{16}{u}\ln(u-1)\right\} + \tau_1^3\tau_2^3\left\{12u - 2u^3 - \frac{8}{u-1} + \frac{8}{u}\ln(u-1)\right\}, \tag{7}$$

where $\tau_{1,2}^3$ denote the third components of the usual isospin operators.

$$S_1(u)^{(\mathrm{III})} = S_1(u)^{(\mathrm{IV})} = (\vec\tau_1\cdot\vec\tau_2 + \tau_1^3\tau_2^3)\left\{\frac{7u^3}{4} - 3u^2 - \frac{u}{2} - 1 + \Big(8u - 3u^3 - \frac{1}{u}\Big)\ln(u-1)\right\}, \tag{8}$$

$$S_1(u)^{(\mathrm{V})} = (\vec\tau_1\cdot\vec\tau_2 + \tau_1^3\tau_2^3)\left\{4u^2 - 3u^3 + 6u - 4 - \frac{4}{u}(u^2-1)^2\ln(u-1)\right\}, \tag{9}$$

$$\begin{aligned} S_1(u)^{(\mathrm{VII})} &= \tau_1^3\tau_2^3\left\{\frac{u^3}{2} + 2u^2 - 11u + 10 + \frac{2}{u}(5u^4 - 14u^2 + 5)\ln(u-1)\right.\\ &\quad \left. + 4\oint_1^{u/2} dx\,\frac{(2-ux)(u^2+2ux-4)}{(u-2x)y}\,\ln\frac{u(x+y)-1}{u(x-y)-1}\right\}, \end{aligned} \tag{10}$$

with the abbreviation $y = \sqrt{x^2-1}$.

$$\begin{aligned} S_1(u)^{(\mathrm{VIII})} &= (\tau_1^3 + \tau_2^3 - \tau_1^3\tau_2^3 - \vec\tau_1\cdot\vec\tau_2)\left\{\frac{1}{2u}(3 + 2u^2 - u^4)\ln(u-1)\right.\\ &\quad \left. + \frac{7u^3}{8} - \frac{5u^2}{2} + \frac{3u}{4} + \frac{3}{2} + u\int_1^{u/2}\frac{dx}{y}(2-ux)\ln\frac{u(x+y)-1}{u(x-y)-1}\right\}, \end{aligned} \tag{11}$$

$$\begin{aligned} S_1(u)^{(\mathrm{X})} &= (\tau_1^3 + \tau_2^3 + \tau_1^3\tau_2^3 + \vec\tau_1\cdot\vec\tau_2)\left\{\frac{1}{2u}(u^4 - 4u^2 - 1)\ln(u-1)\right.\\ &\quad + \frac{u-2}{8}(2 + 2u - u^2) + \frac{1}{4}(2 - 9u^2)\sqrt{u^2-4} + \frac{2}{u}(u^4+1)\ln\frac{u+\sqrt{u^2-4}}{2}\\ &\quad \left. + 2u\oint_1^{u/2}\frac{dx}{u-2x}\left[2uy + (2-u^2)\ln\frac{u-x+y}{u-x-y}\right]\right\}, \end{aligned} \tag{12}$$

$$\begin{aligned} S_1(u)^{(\mathrm{XI})} &= (\tau_1^3 + \tau_2^3 + 3\tau_1^3\tau_2^3 - \vec\tau_1\cdot\vec\tau_2)\left\{2u^2 - \frac{3u^3}{4} - \frac{u}{2} - 1 + \Big(u - \frac{1}{u}\Big)\ln(u-1)\right.\\ &\quad + \frac{1}{4}(2 - 9u^2)\sqrt{u^2-4} + \frac{2}{u}(u^4+1)\ln\frac{u+\sqrt{u^2-4}}{2} + u\int_1^{u/2}\frac{dx}{y}(ux-2)\\ &\quad \left. \times\ln\frac{u(x+y)-1}{u(x-y)-1} + 2u\oint_1^{u/2}\frac{dx}{u-2x}\left[2uy + (2-u^2)\ln\frac{u-x+y}{u-x-y}\right]\right\}, \end{aligned} \tag{13}$$

$$
\begin{aligned}
S_1(u)^{(\mathrm{XII})} &= (\tau_1^3+\tau_2^3+\tau_1^3\tau_2^3+2-\vec{\tau}_1\cdot\vec{\tau}_2)\bigg\{\frac{1}{2u}(4u^2+1-u^4)\ln(u-1) \\
&+\frac{u-2}{8}(u^2-2u-2)+\frac{1}{4}(2-9u^2)\sqrt{u^2-4}+\frac{2}{u}(u^4+1)\ln\frac{u+\sqrt{u^2-4}}{2} \\
&+2u\oint_1^{u/2}\frac{dx}{u-2x}\Big[2uy+(2-u^2)\ln\frac{u-x+y}{u-x-y}\Big]\bigg\}\,, \qquad (14)
\end{aligned}
$$

$$
\begin{aligned}
S_1(u)^{(\mathrm{XIII})} &= (\tau_1^3+\tau_2^3-\tau_1^3\tau_2^3+2+\vec{\tau}_1\cdot\vec{\tau}_2)\bigg\{\frac{1}{2u}(u^4-4u^2-1)\ln(u-1) \\
&+\frac{u-2}{8}(2+2u-u^2)+\frac{1}{4}(2-9u^2)\sqrt{u^2-4}+\frac{2}{u}(u^4+1)\ln\frac{u+\sqrt{u^2-4}}{2} \\
&+2u\oint_1^{u/2}\frac{dx}{u-2x}\Big[2uy+(2-u^2)\ln\frac{u-x+y}{u-x-y}\Big]\bigg\}\,, \qquad (15)
\end{aligned}
$$

$$
S_1(u)^{(\mathrm{XIV})} = (\vec{\tau}_1\cdot\vec{\tau}_2-1)\bigg\{(u^2-2)\sqrt{u^2-4}-\frac{8}{u}\ln\frac{u+\sqrt{u^2-4}}{2}\bigg\}\,, \qquad (16)
$$

$$
\begin{aligned}
S_1(u)^{(\mathrm{XV+XVI})} &= (\vec{\tau}_1\cdot\vec{\tau}_2-\tau_1^3\tau_2^3)\bigg\{4u-8-\frac{8}{u}\ln(u-1) \\
&+(2-3u^2)\sqrt{u^2-4}+\frac{8}{u}(1+u^2)\ln\frac{u+\sqrt{u^2-4}}{2}\bigg\}\,. \qquad (17)
\end{aligned}
$$

The contribution from diagram (XV) together with the irreducible part of diagram (XVI) is proportional to the difference of their isospin factors on the left nucleon line: $[1+\tau_1^3,\tau_1^c]\epsilon^{abc} = 2i(\tau_1^a\delta^{b3}-\tau_1^b\delta^{a3})$. Furthermore, a closer inspection of the spin-dependent factors reveals that the longitudinal spectral function $S_2(u)$ receives only contributions from those diagrams where the photon does not couple to a pion in flight. The few possible contributions to $S_2(u)$ from the diagrams (II), (XIV) and (XV+XVI) read for $u\geq 2$:

$$
\begin{aligned}
S_2(u)^{(\mathrm{II})} &= \vec{\tau}_1\cdot\vec{\tau}_2\bigg\{\frac{u^3}{3}-\frac{5u}{2}+\frac{10}{3}-\frac{1}{u}-\frac{2}{u^2}-\frac{2}{u^3}(u^2-1)^2\ln(u-1)\bigg\} \\
&+\tau_1^3\tau_2^3\bigg\{\frac{5u}{2}-2-\frac{3}{u}-\frac{6}{u^2}-\frac{2}{u^3}(u^4+3)\ln(u-1)\bigg\}\,, \qquad (18)
\end{aligned}
$$

$$
S_2(u)^{(\mathrm{XIV})} = (1-\vec{\tau}_1\cdot\vec{\tau}_2)\bigg\{\frac{u^2-1}{6u^2}(u^2+6)\sqrt{u^2-4}+\frac{2}{u^3}(2u^2-2-u^4)\ln\frac{u+\sqrt{u^2-4}}{2}\bigg\}\,, \qquad (19)
$$

$$
\begin{aligned}
S_2(u)^{(\mathrm{XV+XVI})} &= (\vec{\tau}_1\cdot\vec{\tau}_2-\tau_1^3\tau_2^3)\bigg\{\frac{2}{u^2}+\frac{1}{u}-\frac{10}{3}+\frac{5u}{2}-\frac{u^3}{3}+\frac{2}{u^3}(u^2-1)^2\ln(u-1) \\
&+\frac{u^2-1}{6u^2}(u^2+6)\sqrt{u^2-4}+\frac{2}{u^3}(2u^2-2-u^4)\ln\frac{u+\sqrt{u^2-4}}{2}\bigg\}\,. \qquad (20)
\end{aligned}
$$

We have also checked gauge-invariance. The (total) spectral functions $S_1(u)$ and $S_2(u)$ stay ξ-independent when adding a longitudinal part to the photon propagator: $g_{\mu\nu}\to g_{\mu\nu}+\xi\,k_\mu k_\nu$. The "encircled" integrals appearing in Eqs.(10,12–15) symbolize the following regularization prescription:

$$
\oint_1^{u/2}dx\,\frac{f(x)}{u-2x} = \int_1^{u/2}dx\,\frac{f(x)-f(u/2)}{u-2x}\,. \qquad (21)
$$

This regularization prescription eliminates from the contributions of some diagrams ((V), (VII), (X), (XI), (XII), (XIII)) to the spectral function $S_1(u)$ an infrared singularity arising from the emission of soft photons ($\bar{N}N \to \pi\pi\gamma_{\rm soft}$). The singular factor $(u-2x)^{-1}$ originates from a pion propagator. The regularization prescription defined in Eq.(21) is equivalent to the familiar "plus"-prescription employed commonly for parton splitting functions in order to eliminate there an analogous infrared singularity due to soft gluon radiation [13]. It has also been used in our previous work in ref.[5]. We note as an aside that the non-elementary integrals ($\int_1^{u/2} dx \dots$) in Eqs.(10–15) can be solved in terms of dilogarithms and squared logarithms.

With the help of the spectral functions ${\rm Im}V_{S,T}(i\mu)$ or $S_{1,2}(u)$ the $2\pi\gamma$-exchange spin-spin and tensor potentials in coordinate space $\widetilde V_{S,T}^{(2\pi\gamma)}(r)$ can be calculated easily via a modified Laplace transformation:

$$\begin{aligned}\widetilde V_S^{(2\pi\gamma)}(r) &= \frac{1}{6\pi^2 r}\int_{2m_\pi}^\infty d\mu\, \mu e^{-\mu r}\{\mu^2\, {\rm Im}V_T(i\mu) - 3\,{\rm Im}V_S(i\mu)\} \\ &= \widetilde V_S^{(0)}(r) + \vec\tau_1\cdot\vec\tau_2\, \widetilde W_S^{(0)}(r) + \tau_1^3\tau_2^3\, \widetilde V_S^{(\rm cib)}(r) + (\tau_1^3+\tau_2^3)\, \widetilde V_S^{(\rm csb)}(r)\,, \end{aligned} \tag{22}$$

$$\begin{aligned}\widetilde V_T^{(2\pi\gamma)}(r) &= \frac{1}{6\pi^2 r^3}\int_{2m_\pi}^\infty d\mu\, \mu e^{-\mu r}\,(3+3\mu r+\mu^2 r^2)\, {\rm Im}V_T(i\mu) \\ &= \widetilde V_T^{(0)}(r) + \vec\tau_1\cdot\vec\tau_2\, \widetilde W_T^{(0)}(r) + \tau_1^3\tau_2^3\, \widetilde V_T^{(\rm cib)}(r) + (\tau_1^3+\tau_2^3)\, \widetilde V_T^{(\rm csb)}(r)\,, \end{aligned} \tag{23}$$

where we have given in the second line of Eqs.(22,23) the decomposition into isospin-conserving (0), charge-independence breaking (cib), and charge-symmetry breaking (csb) parts. The numbers in the second to fifth row of Tables I and II display the dropping of the spin-spin and tensor potentials $\widetilde V_{S,T}^{(0)}(r)$, $\widetilde W_{S,T}^{(0)}(r)$, $\widetilde V_{S,T}^{(\rm cib)}(r)$ and $\widetilde V_{S,T}^{(\rm csb)}(r)$ with the nucleon distance r in the interval $1.0\,{\rm fm} \leq r \leq 1.9\,{\rm fm}$. One notices quite sizeable charge-independence breaking (cib) spin-spin and tensor potentials. Their magnitudes (for example, $-0.11\,$MeV and $0.09\,$MeV at a nucleon distance of $r = m_\pi^{-1} = 1.4\,$fm) are exceptionally large in comparison to all so far known isospin-breaking spin-spin and tensor potentials generated by long-range pion-exchange. In particular, the potentials $\widetilde V_{S,T}^{(\rm cib)}(r)$ exceed their counterparts from (leading and next-to-leading order) one-loop $\pi\gamma$-exchange by more than one order of magnitude (see herefore Table I in ref.[6]). As can be seen from the last row in Tables I and II the next largest $2\pi\gamma$-exchange potentials are the charge-symmetry breaking (csb) ones. The spin-spin and tensor potentials $\widetilde V_{S,T}^{(\rm csb)}(r)$ reach about 40% of the charge-independence breaking ones and they have throughout the same sign as $\widetilde V_{S,T}^{(\rm cib)}(r)$. We note also that in each case the numerically dominant contributions are coming from the diagrams (V), (VII), (X)–(XIII). Chiral power counting [4] is obviously respected in our calculation since all contributions to the spectral functions ${\rm Im}V_{S,T}(i\mu)$ scale in the same way with fine-structure constant α and the small mass and momentum scales m_π and μ. The sizable values of the isospin-violating spin-spin and tensor potentials $\widetilde V_{S,T}^{(\rm cib,csb)}(r)$ are due to the large value of the low-energy constant $c_4 = 3.4\,{\rm GeV}^{-1}$ and the large number of contributing diagrams.

In summary, we have calculated in this work the electromagnetic corrections to the two-pion exchange nucleon-nucleon interaction that is generated by the isovector spin-flip $\pi\pi NN$ contact-vertex proportional to the large low-energy constant $c_4 = 3.4\,{\rm GeV}^{-1}$. Our analytical results for these novel and exceptionally large isospin-violating spin-spin and tensor potentials have been presented in a form such that they can be easily implemented into future NN-phase shift analyses and few-body calculations.

r [fm]	1.0	1.1	1.2	1.3	1.4	1.5	1.6	1.7	1.8	1.9
$\widetilde{W}_S^{(2\pi)}$ [MeV]	24.95	13.61	7.78	4.63	2.84	1.80	1.16	0.769	0.518	0.354
$\widetilde{V}_S^{(0)}$ [keV]	–545	–273	–144	–79.3	–45.4	–26.8	–16.3	–10.1	–6.44	–4.17
$\widetilde{W}_S^{(0)}$ [keV]	240	126	68.9	39.5	23.4	14.3	8.97	5.76	3.77	2.51
$\widetilde{V}_S^{(\mathrm{cib})}$ [keV]	–1307	–660	–351	–195	–112	–66.8	–40.9	–25.6	–16.4	–10.7
$\widetilde{V}_S^{(\mathrm{csb})}$ [keV]	–559	–280	–148	–81.9	–47.0	–27.8	–16.9	–10.6	–6.72	–4.37

Table I: The dominant isovector 2π-exchange spin-spin potential $\widetilde{W}_S^{(2\pi)}(r)$, and electromagnetic corrections to it, as a function of the nucleon distance r. The values in the fourth and fifth row correspond to the isospin-violating spin-spin potentials $\widetilde{V}_S^{(\mathrm{cib})}(r)$ and $\widetilde{V}_S^{(\mathrm{csb})}(r)$.

r [fm]	1.0	1.1	1.2	1.3	1.4	1.5	1.6	1.7	1.8	1.9
$\widetilde{W}_T^{(2\pi)}$ [MeV]	–22.91	–12.35	–6.98	–4.10	–2.49	–1.56	–0.996	–0.652	–0.434	–0.294
$\widetilde{V}_T^{(0)}$ [keV]	452	224	117	63.9	36.2	21.2	12.8	7.89	4.98	3.20
$\widetilde{W}_T^{(0)}$ [keV]	–210	–108	–59.0	–33.4	–19.6	–11.9	–7.38	–4.69	–3.05	–2.01
$\widetilde{V}_T^{(\mathrm{cib})}$ [keV]	1098	549	289	159	91.0	53.7	32.5	20.2	12.8	8.31
$\widetilde{V}_T^{(\mathrm{csb})}$ [keV]	469	233	122	66.7	37.9	22.3	13.4	8.31	5.26	3.39

Table II: The dominant isovector 2π-exchange tensor potential $\widetilde{W}_T^{(2\pi)}(r)$, and electromagnetic corrections to it, as a function of the nucleon distance r. The values in the fourth and fifth row correspond to the isospin-violating tensor potentials $\widetilde{V}_T^{(\mathrm{cib})}(r)$ and $\widetilde{V}_T^{(\mathrm{csb})}(r)$.

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